Spatiotemporal evolution of high-power relativistic laser pulses in electron-positron-ion plasmas

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The spatiotemporal pulse dynamics of a high-power relativistic laser pulse interacting with an electron-positron-ion plasmas is investigated theoretically and numerically. The occurrence of pulse compression is studied. The dependence of the mechanism on the concentration of the background ions in electron positron plasma is emphasized.

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I. INTRODUCTION

At high laser intensities the nonlinear interaction between the laser beam and the surrounding plasma medium becomes important, giving rise to a variety of physical effects [1–4], which are not observed in the linear regime. These include relativistic optical guiding, harmonic excitation, wake-field generation, laser pulse frequency shifting, and pulse compression, to mention a few. Akhiezer and Polovin [1] investigated analytically the propagation of very intense electromagnetic radiation in overdense plasmas and provided exact traveling wave solutions. Dawson [2] investigated nonlinear longitudinal electron oscillations in an infinite field-free cold plasma. Sprangle *et al.* [3] developed a one-dimensional (1D) nonlinear theory to model plasma wake-field generation, relativistic optical guiding and coherent harmonic radiation production in intense laser plasma interactions. Pukhov [4] discussed in his elegant review some of the important physical effects emerging at relativistic laser intensities. In earlier work [5], we have also put forward the possibility of laser pulse amplification via a colliding beam scheme in a plasma medium.

The interaction of ultraintense very short laser pulses with plasmas [6–11] has attracted a great deal of attention recently, both in fundamental research and for technological applications, such as particle acceleration, inertial confinement fusion, high harmonic generation, and x-ray lasers [12–18]. The standard approach to produce ultrashort, ultraintense multiterawatt laser pulse is the chirped-pulse-amplification (CPA) technique [19], in which a laser pulse is stretched, amplified and recompressed. The CPA scheme has shown the ability of generating subpicosecond petawatt laser pulses with up to 500 Joules per pulse. This approach is limited by the finite bandwidth of the active *mm* amplifiers used in lasers. Ross *et al.* [20] investigated a new scheme of parametric amplification to produce ultrashort and powerful pulses. Super-radiant amplification of an ultrashort laser pulse was observed by Shvets *et al.* [21] who considered an electromagnetic (em) beam colliding with a long counter-propagating low-intensity pump in the plasma. The methods reported in Refs. [20,21]. need at least two counterpropagating laser pulses, which makes practical realization of those methods difficult.

Relativistic mass variation during laser-plasma interaction is the origin for longitudinal self-compression of a laser pulse down to a single laser cycle in length, with a corresponding increase in intensity. The main source of nonlinearity is the relativistic mass increase due to the quiver motion of the electrons in the field of the laser. In the last few years, several scenarios have been proposed for the self-focusing and self-compression of a laser pulse in plasma [22]. Shorokhov *et al.* [23] have employed a three-dimensional (3D) particle-in-cell (PIC) simulation to show that a 30 fs long laser pulse is efficiently compressed to 5 fs by using a periodic plasma-vacuum structure to damp filamentation. Tsung *et al.* [24] reported a scheme to generate single-cycle laser pulses based on photon deceleration in underdense plasmas. This robust and tunable process is ideally suited for lasers above critical power because it takes advantage of the relativistic self-focusing of these lasers and the nonlinear features of the plasma wake. Ren *et al.* [25] demonstrated the compression and focusing of a short laser pulse by a thin plasma lens. A set of analytical formulas for the spot size and for the length evolution of a short laser pulse were derived in their model. Shibu *et al.* [26] also proposed the possibility of pulse compression in relativistic homogeneous plasma and reported the interplay between transverse focusing and longitudinal
compression.

It was recently demonstrated numerically [27] that microwave pulses can be compressed (with or without a frequency shift) inside a magnetized plasma by changing the magnitude or the direction of the magnetic field uniformly in space and adiabatically in time. Balakin *et al.* [28] have investigated the self focusing of few optical cycle pulses recently. They showed that the wave-field self-focusing proceeds with overtaking the steepening of the pulse longitudinal profile, leading to shock-wave formation. Consequently, a more complex singularity is formed where an unlimited field increase is followed by wave breaking with a broad power-law pulse spectrum.

Most of earlier work outlined above [23–27] has focused on pulse self-compression in homogeneous relativistic plasma in relation with short ultraintense pulses. Some studies do exist on laser pulse propagation in inhomogeneous (nonuniform) plasmas [29–32], but these are limited to transverse beam focusing. Longitudinal pulse compression needs more attention to be paid to, both analytically and numerically, due to its relevance in the development of a practical model high-power ultrashort laser. A number of investigations have focused on the propagation of an em beam in axially inhomogeneous plasma, addressing effects such as the relativistic self focusing of intense laser radiation in axially inhomogeneous plasma in relation with the fast ignitor scheme [29], the propagation regime of circularly polarized beam in relativistic inhomogeneous plasma with an indication of em beam penetration in overdense plasma [30], the study of intense laser beam propagation in axially inhomogeneous relativistic plasma [31], the cumulative effects of an axially inhomogeneous plasma column on the self focusing and third harmonic generation of an em beam when the scale length of the inhomogeneity is larger than the wavelength [32].

Modern lasers (e.g., Vulcan Petawatt Upgrade at Rutherford Appleton Laboratory's Central Laser Facility and the Gekko Petawatt Laser at the Gekko XII facility in the Institute of Laser Engineering at Osaka University) rely on CPA techniques for amplifying an ultrashort laser pulse to extremely large intensities. Because of limitations e.g., in gain bandwidth, high-power CPA systems are currently limited from below to pulses of order 30 fs. The physical reason for this limitation is the finite bandwidth of the active medium amplifiers used in the lasers. The advantage of plasma as an "active" medium for pulse compression is that it sustains extremely high intensities. Nonlinearity becomes significant only close to the relativistic threshold and thus high power can be achieved.

In this paper, we are interested in investigating the self-compression and self-focusing of a Gaussian-cross-section relativistic laser pulse propagating in an electron-positron-ion plasma. An earlier investigation by Berezhiani and Mahajan [33] has focused on the nonlinear interaction of circularly polarized electromagnetic waves in unmagnetized cold *e-p-i* plasma, in fact exploring the possibility of finding a 1D soliton solution in such a plasma configuration. In a recent research Mahajan and Shatashvili [34] investigated the nonlinear propagation of high intensity em waves in a pair ion plasma contaminated with a small fraction of a high mass immobile ions (for symmetry breaking). They highlighted a very remarkable property of this new and interesting state of (laboratory created) matter, that it can strongly localize the em radiation with finite density excess in the region of localization. Another recent approach by Mahajan *et al.* [35] demonstrated the nonlinear propagation of electromagnetic waves in pair plasmas, in which the electrostatic potential plays a very important role of a binding glue. It is shown that the temperature asymmetry leads to a (localizing) nonlinearity that is qualitatively different from the ones originating from the intrinsic mass or density difference. The temperature-asymmetry-driven focusing-defocusing nonlinearity supports stable localized wave structures in 13 dimensions, which, for certain parameters, may have flat-top shapes.

However, we focus here on the spatiotemporal dynamics of the pulse, an aspect not covered earlier. We rely on a nonlinear Schrödinger equation (NLSE) to study the spatiotemporal dynamics of the em field envelope. Following [36], we introduce a set of trial functions via the intensity profile of the laser pulse, and follow their evolution in space/time in the plasma. At a first step, we adopt a 1D model, relying on the em wave equation as derived from Maxwell's equations. A nonlinear Schrodinger equation is then obtained and solved by using the paraxial formalism, to model the occurrence of longitudinal pulse width compression and associated energy localization. The analysis is then extended to a 3D pulse profile description. A pair of appropriate trial functions are defined, accounting for the beam width (in space) and the pulse duration (in time), whose evolution describes the dynamics of the pulse. Both longitudinal and transverse self compression are examined for a finite extent Gaussian laser pulse through this model. These functions are determined by a system of coupled nonlinear differential equations, which are integrated numerically to yield the spatiotemporal laser pulse profile.

II. ANALYTICAL MODEL

We consider a three-component *e-p-i* plasma consisting of electrons and positrons (denoted by the minus/plus subscript, respectively, everywhere below) of opposite charge $q_+=-q_-=e$ and equal masses $m_\mp=m$, in addition to a massive ion component in the background. Because of their large inertia, the ions do not respond to the dynamics under consideration and just provide a neutralizing background, hence we take their density to be $n_i=n_{i,0}=$constant. The equilibrium state is characterized by an overall charge neutrality $n_{-,0}=n_{+,0}+n_i$, where $n_{\mp,0}$ are the unperturbed number densities of the electrons/positrons (denoted by the minus/plus subscript respectively).

In order to describe the propagation of electromagnetic waves in such a plasma, we may use as a starting point Maxwell's equations. The system's evolution in terms of the vector ($\mathbf{A}$) and the scalar (ϕ) potentials is governed by,

$$\mathbf{E}=-\frac{1}{c}\frac{\partial \mathbf{A}}{\partial t}-\nabla\phi, \quad \mathbf{B}=\nabla\times\mathbf{A}. \tag{1}$$

In rescaled form, the field equations are expressed as [33]:

$$\frac{\partial^2 A}{\partial t^2}-\Delta\mathbf{A}+\frac{\partial}{\partial t}\nabla\,\phi+[n_-v_- -(1-\alpha)n_+v_+]=0, \tag{2}$$

and

$$\Delta\phi=[n_- -(1-\alpha)n_+ -\alpha]. \tag{3}$$

The system is closed by invoking the hydrodynamic equations, consisting of the equation of motion

$$\frac{\partial P_\pm}{\partial t}+\nabla[1+(P_\pm)^2]^{1/2}=\mp\frac{\partial \mathbf{A}}{\partial t}\mp\nabla\phi, \tag{4}$$

and the continuity equation

$$\frac{\partial n_\pm}{\partial t}+\nabla(n_\pm v_\pm)=0, \tag{5}$$

for each of the mobile components ($\mp$). Equations (2)–(5), expressed in the gauge $\nabla\cdot\mathbf{A}=0$, are dimensionless with the following normalizations: the time and space variables are measured in units of the inverse electron plasma frequency $\omega_e^{-1}[=(4\pi n_0^- e^2/m_e)^{-1/2}]$, and the collisionless skin depth c/ω_e; the vector and scalar potentials are normalized to

$m_e c^2/e$; the relativistic momentum P to $m_e c$; finally, n^- and n^+ are scaled by their respective equilibrium densities $n_{-,0}$ and $n_{+,0}$. The coefficient $\alpha = n_i/n_{-,0}$ denotes the ratio of the ion density to the electron equilibrium density.

We consider the propagation of circularly polarized em pulse along the z direction. The vector potential for em pulse can be written as:

$$A(r,z,t) = A(r,z,t)(e_x + ie_y)\exp[i(kz - \omega t)], \quad (6)$$

where ω and k are, respectively, the frequency and wave number of pulse. The wave number of the em beam satisfies the plasma dispersion relation, $c^2k^2 = \omega^2 - \omega_e^2 - \omega_p^2$, where ω_e (ω_p) is the electron (positron) plasma frequency. An elegant kinetic description of various linear modes in a nonrelativistic pair magnetoplasma has been investigated by Iwamoto [37]. Zank and Greaves [38] discussed the linear properties of various electrostatic and electromagnetic modes in unmagnetized and in magnetized pair plasmas, and also considered two-stream instability and nonenvelope (pulse) solitary wave solutions.

The propagation of an em pulse in e-p-i plasma is described by the following set of equations

$$2ik\frac{\partial A}{\partial z} + \frac{2-\alpha}{v_g^2\omega^2}\frac{\partial^2 A}{\partial \tau^2} + \left(\frac{\partial^2 A}{\partial r^2} + \frac{1}{r}\frac{\partial A}{\partial r}\right) + A\frac{\phi}{(1-\phi^2)}[\alpha - (2-\alpha)\phi] = 0, \quad (7)$$

and

$$\frac{\partial^2\phi}{\partial z^2} = \frac{1}{2}\left[\frac{(1+|A|^2)}{(1+\phi)^2} - \frac{(1-\alpha)(1+|A|^2)}{(1-\phi)^2} - \alpha\right], \quad (8)$$

[33], where $\tau = t - (z/v_g)$ and v_g is the group velocity of em pulse. Assuming that the (normalized) spatial extension (length) of the pulse is large, i.e., satisfies $L_\parallel \gg (1+|A|^2)^{1/2}$, then Eq. (8) implies

$$\phi = \frac{1}{2}\frac{\alpha}{2-\alpha}\frac{|A|^2}{(1+|A|^2)}. \quad (9)$$

It is evident from the above expression that ϕ is proportional to α (for $\alpha \ll 1$), and in fact vanishes in the absence of background ions. It was shown in [39] that a pure electron positron plasma ($\alpha=0$) cannot sustain an electrostatic potential (ϕ). As a result, a circularly polarized electromagnetic wave cannot be localized in pure electron positron plasma. Also, $\phi \ll 1$ even for a large electromagnetic pulse. Combining Eqs. (9) and (7) we are led to the nonlinear Schrödinger Equation (NLSE)

$$2ik\frac{\partial A}{\partial z} + \frac{2-\alpha}{v_g^2\omega^2}\frac{\partial^2 A}{\partial \tau^2} + \left(\frac{\partial^2 A}{\partial r^2} + \frac{1}{r}\frac{\partial A}{\partial r}\right) + \frac{\alpha^2}{4(2-\alpha)}A[1-(1+|A|^2)^{-2}] = 0, \quad (10)$$

where we have neglected terms of cubic or higher order in ϕ. The second, third (within parentheses) and fourth contributions to the left-hand side represent diffraction, dispersion and nonlinearity. Note that the last (nonlinear) term becomes $\simeq \frac{\alpha}{2}|A|^2A$ in the small $|A|$ limit, recovering the cubic NLS derivation.

Equation (10) is the main outcome of this section, which will now be relied upon, to proceed by studying the beam profile dynamics. Note that nonlinearity is entirely due to the ion background. We stress that have derived Eq. (10) by considering the longitudinal and transverse spatial dynamics of an em beam propagating in e-p-i plasma system. This is clearly distinct from the derivation of Berezhiani and Mahajan [33], who have derived a similar-structured equation for the 1D dynamics of a circularly polarized em beam, and also clearly different from the equation(s) derived in various forms and via different physical approaches in earlier works [22,36].

III. BEAM PROFILE DYNAMICS

To proceed, we shall follow the paraxial beam profile approach, as introduced by Sharma and co-workers in Ref. [36]. We consider a beam whose initial profile presents a Gaussian intensity distribution, viz. $A^2(r,z=0,t) = A_0^2\exp(-r^2/r_0^2)\exp(-t^2/\tau_0^2)$, where $r = (x^2+y^2)^{1/2}$ is the radial coordinate (in cylindrical polar coordinates), τ_0 is the initial pulse width (in time) and r_0 is the initial spot size (in space) of the pulse. The solution of Eq. (10) can be expressed as [36]

$$A(r,z,t) = A_0(r,z,t)\exp[-ikS(r,z,t)]. \quad (11)$$

where both amplitude (A_0) and eikonal (S) are real quantities. The eikonal S is related with the curvature of wavefront, while the amplitude (square) represents the intensity profile. Substituting for A from Eq. (11) in Eq. (10) and separating the real from the imaginary parts, one obtains

$$\frac{\partial A_0^2}{\partial z} + \frac{\partial S}{\partial r}\frac{\partial A_0^2}{\partial r} + A_0^2\left(\frac{\partial^2 S}{\partial r^2} + \frac{1}{r}\frac{\partial S}{\partial r}\right)$$

$$\frac{\partial S}{\partial t}\frac{\partial (A_0^2)}{\partial t} + A_0^2\left(\frac{\partial^2 S}{\partial t^2}\right) = 0, \quad (12)$$

and

$$2\frac{\partial S}{\partial z} + \left(\frac{\partial S}{\partial r}\right)^2 + \left(\frac{\partial S}{\partial t}\right)^2 = \frac{\omega^2\alpha' A[1-(1+|A|^2)^{-2}]}{c^2k^2} + \frac{1}{k^2A_0}\left(\frac{\partial^2 A_0}{\partial r^2} + \frac{1}{r}\frac{\partial A_0}{\partial r}\right) + \frac{1}{k^2A_0}\left(\frac{\partial^2 A_0}{\partial t^2}\right), \quad (13)$$

where $\alpha' = \frac{\alpha^2}{4(2-\alpha)}$. Adopting the paraxial theory [36,40,41], we anticipate a solution for Eqs. (13) and Eq. (12) in the form

$$A_0^2(r,z,t) = \frac{A_{00}^2}{g(z)f(z)^2}\exp\left[\frac{-r^2}{r_0^2f(z)^2}\right]\exp\left[\frac{-t^2}{\tau_0^2g(z)^2}\right], \quad (14)$$

and

$$S(r,z,\tau) = \frac{r^2}{2f}\frac{df}{dz} + \frac{t^2}{2g}\frac{dg}{dz} + \phi(z), \quad (15)$$

where $f(z)$ and $g(z)$ are beam width (space) and pulse width (time) parameters, whose evolution [governed by Eqs. (17) and (18), see below] determines the beam dynamics. Identifying the components of the eikonal (S) in the latter expression, the first term above is indicative of the spherical curvature of the wavefront, while ϕ represents its departure from the spherical nature. The latter two Eqs. (14) and (15) have also recently been employed in [36] for Gaussian beam propagation in collisional plasmas.

Using Eqs. (11), (14), and (15), the spatiotemporal evolution of the em pulse in *e-p-i* plasma can be written as,

$$A^2(z,\tau)=\frac{A_0^2[T(0)R(0)^2]}{T(z)R(z)^2}\exp\left[-\frac{\tau^2}{T(z)^2}-i\frac{\epsilon_0^{1/2}}{2}\frac{\tau^2}{T(\zeta)}\frac{dT}{dz}-\frac{\rho^2}{R(z)^2}-i\frac{\epsilon_0^{1/2}}{2}\frac{\rho^2}{R(z)^2}\frac{d\rho}{dz}\right], \quad (16)$$

where $T(z)=\tau_1 z$ is the pulse width (time) in plasma and $R(z)=\rho_1 z$ is the beam width (space) in the radial direction in plasma. Here, $T(0)=\tau_1=\tau_0\omega$ is the initial dimensionless pulse width (at $t=0$), and $R(0)=\rho_1=\rho_0\omega/c$ is the initial dimensionless beam width.

The EM pulse profile in plasma can be obtained by solving the following two coupled second order ODEs for the self-focusing parameter $f(z)$ and self-compression parameter $g(z)$:

$$\epsilon_0\frac{d^2f}{d\xi^2}=\frac{1}{\rho_1^4 f^3}-\frac{(2\alpha')A_0^2}{\rho_1^2 f^3 g(1+A_0^2/f^2g)^3}, \quad (17)$$

$$\epsilon_0\frac{d^2g}{d\xi^2}=\frac{1}{\tau_1^4 g^3}-\frac{(2\alpha')A_0^2}{\tau_1^2 f^2 g^2(1+A_0^2/f^2g)^3}. \quad (18)$$

For an initial plane wave the boundary conditions on Eqs. (17) and (18) are taken at $\xi=0$ as

$$\frac{df}{d\xi}=\frac{dg}{d\xi}=0 \quad \text{and} \quad f=g=1.$$

Equations (17) and (18) can be numerically integrated using the initial boundary conditions above, to evaluate the beam width parameter f and pulse width parameter g as a function of z. The numerical estimation of f and g as a function of propagation distance will allow us to predict the variation of pulse width and beam radius in *e-p-i* plasma.

Concluding this part, we have derived a set of coupled envelope equations for the evolution of the spot size and the pulse (time) length. This clearly illustrates the simultaneously operating processes of focusing and compression in a three-component plasma. These coupled equations enable us to consider now how self-focusing may affect (and in fact accelerates) the pulse compression rate. We emphasize that these equations are idealized, since we have assumed that one spot size describes the beam at any instant t.

IV. NUMERICAL INVESTIGATION

We have performed an extensive numerical investigation of the beam profile dynamics for the following laser plasma parameters [25]: $I_0=1.37\times10^{18}$ W/cm^2, $\lambda=1$ μm, $r_0=20$ μm, $\tau_0=10$ psec, $n_0=4\times10^{20}$ cm^3, $\omega=10^{15}$ rad sec^{-1}, and $\alpha=0.01$. We have numerically integrated Eqs. (17) and (18) (using the initial boundary conditions stated above), to evaluate the beam width parameter f and pulse width parameter g as functions of the distance of propagation z. The numerical evaluation of f and g is used to obtain the spa-

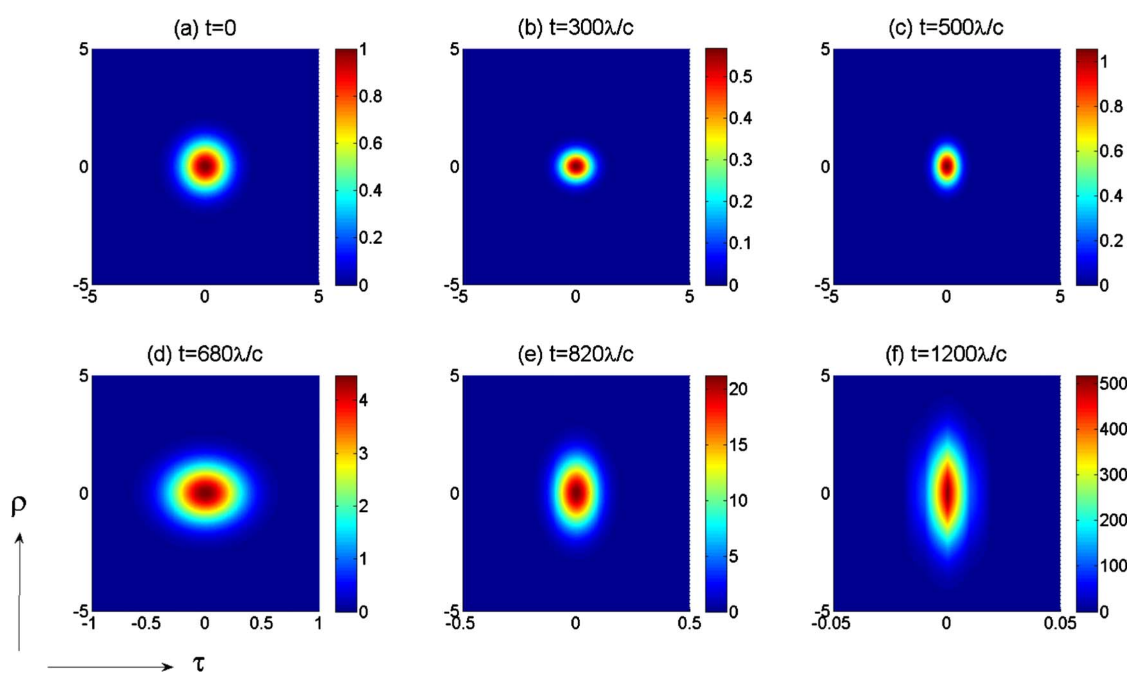

FIG. 1. (Color online) Normalized intensity ($\mathbf{A}^2$) plots as a function of time. The side-bar shows the variation in normalized intensity. The region near to center (*ρ=0, τ=0*) shows maximum intensity (corresponding to the top of bar) while as we move away from the central (*ρ=0, τ=0*) region, the intensity decreases (corresponding to the middle and lower part of the bar respectively).

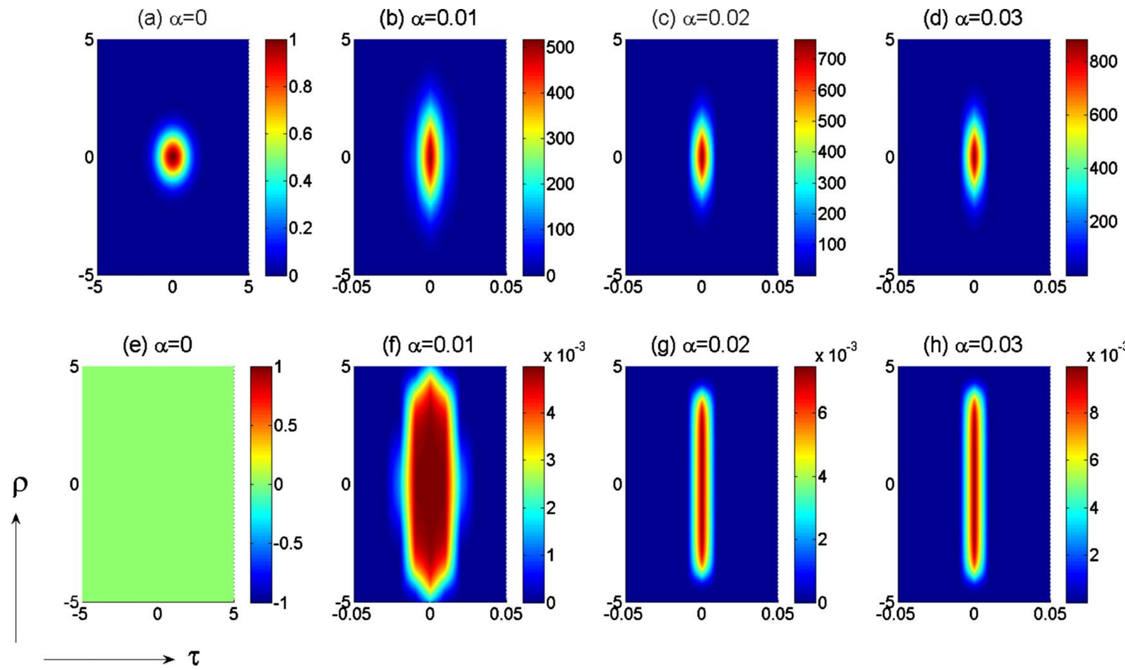

FIG. 2. (Color online) The dependence of normalized intensity ($\mathbf{A}^2$) plots (upper row) and potential (ϕ) plots (lower row) as a function of the ion-to-electron density ratio (α). The side-bar shows the variation in normalized intensity (in upper row) and potential (in lower row). The region near to center ($\rho=0$, $\tau=0$) shows maximum intensity/potential (corresponding to the top of bar) while as we move away from the central ($\rho=0$, $\tau=0$) region, the intensity/potential decreases (corresponding to the middle and lower part of the bar respectively).

tiotemporal beam intensity profile [as given by Eq. (16)] in a three-component plasma at different time instants t (equivalent to a fixed distance z).

In Fig. 1, the intensity profile of a Gaussian beam is depicted at different time t=0, 1, 3, 9, 12, and 14.5 (in units of *psec*). The plots in Fig. 1 depict the transverse focusing of the laser pulse, which is followed by a longitudinal pulse compression due to the combined effect of relativistic mass variation and admixture of ions in electron-positron plasma. The numerical results suggest that transverse focusing competes with the process of longitudinal self compression.

Equation (9) shows the proportionality of scalar potential to α i.e., the ion equilibrium density. It was shown in [39] that a pure electron positron plasma ($\alpha=0$) cannot sustain an electrostatic potential (ϕ). As a result, a circularly polarized electromagnetic wave cannot be localized in pure electron positron plasma. It is now straightforward, to combining Eqs. (16)–(18), as a basis in order to find out the possibility of pulse compression in three component unmagnetized plasma. We will show below how longitudinal pulse compression occurs, in fact sheerly due to the presence of even a small fraction of background ions ($\alpha \ll 1$) in an electron-positron plasma.

An initially Gaussian pulse, with the pulse duration τ_0 =10 psec, beam radius r_0=20 μm, and the amplitude A_0 =1, propagates through a *e-p-i* plasma of density n_0=4 $\times 10^{20}$ cm^3 and α=0, 0.01, 0.02, and 0.03. Figure 2 shows how the pulse compression depends on the presence of α and pulse intensity changes with increase of it, during its propagation in *e-p-i* plasma in the pulse comoving reference system. We observe a compression by more than 5 times for α=0.03, and this process is energetically efficient.

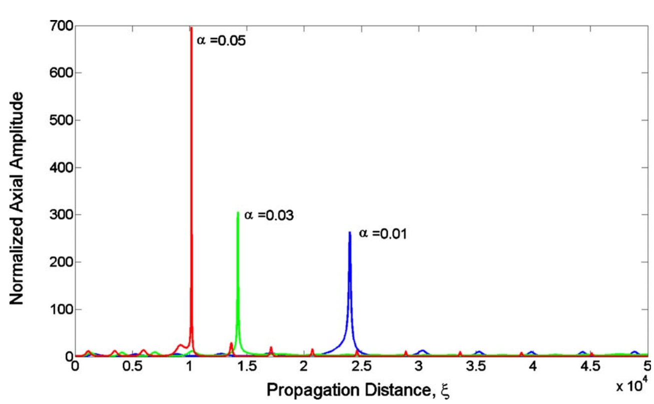

FIG. 3. (Color online) Normalized axial amplitude as a function of propagation distance.

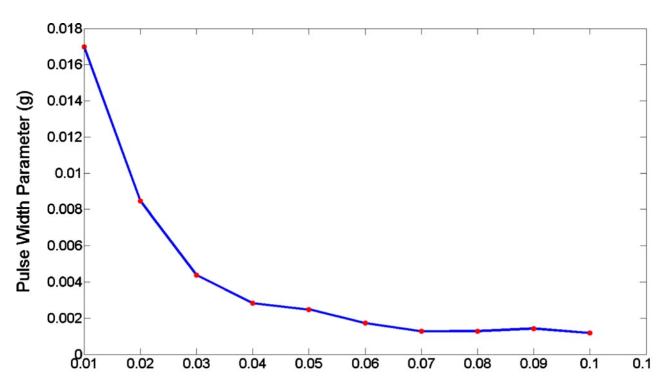

FIG. 4. (Color online) Pulse width parameter (g) as a function of α.

In order to trace the dependence of the normalized axial amplitude on the background ion concentration via α, we have numerically evaluated the axial variation amplitude versus the propagation distance ξ, for difference values of the ion density (expressed via α); the result is shown in Fig. 3. A Gaussian pulse, with initial pulse duration $\tau_0=10$ psec and initial amplitude $A_0=1$, is seen to undergo a periodic variation of its axial amplitude. The same phenomenon is observed for $\alpha=0.01$, 0.03 and 0.05. Figure 3 shows that the successive intensity peaks are higher for higher α, and that these peaks correspond to the foci (beam intensity maxima) corresponding to higher α value. One may conclude that the successive foci become smaller in beam width and pulse width as the ion-to-electron density ratio α decreases. We clearly observe the compression of a picosecond pulse down to the femtosecond range.

We have also depicted the pulse width parameter g for different values of α, for the same parameters as in Fig. 2; the result is shown in Fig. 4. The pulse width parameter is a measure of the longitudinal compression of the initial pulse width $[T(z)=\tau_0\times g(z)]$. We observe from Fig. 4 that a pulse of initial pulse duration 10 psec can be compressed in the range of femtosecond or subattosecond range for $\alpha=0.05$.

V. CONCLUSIONS

From our models, we can identify the mechanism behind the self-compression as a spatiotemporal interplay of plasma-induced refraction, diffraction of the laser beam, and the presence of ions in electron-positron plasmas. The role of *e-p-i* plasma is important in the compression process because it is responsible for the spatiotemporal coupling at high (i.e., ionizing) laser intensities. The numerical results strongly predict that the spatiotemporal dynamics of electromagnetic beam is sensitive due to the presence of ions in e-p plasma system.

In conclusion, we demonstrate a numerical mechanism for spatiotemporal pulse evolution in a plasma of electrons, positrons and massive ions for the first time. Previous investigations focused the em pulse compression in a plasma medium consisting of electrons and massive ions. By propagating intense, femtosecond pulses inside a *e-p-i* plasma system, we demonstrate the temporal self-compression of a pulse with transverse self-focusing. We have explored the possibility of pulse compression in the range of subattosecond. Thus, our work provides impetus for further development of 3D simulation modeling of laser-plasma interactions, demonstrating that fundamentally new physical phenomena can be found. This result represents a significant simplification over other pulse compression techniques at relativistic laser intensities, and will be useful in many applications in high field and plasma science. This mechanism is likely to scale to higher intensities and pulse energies, and of relevance to the use of plasma-based waveguides.

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